

# Particle production and dissipation caused by the Kaluza-Klein tower

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Two-step dissipation is studied in supersymmetric models in which the field in motion couples to bulk fields in the higher dimensional space. Since the Kaluza-Klein tower of the intermediate field changes its mass-spectrum during the evolution, there could be back-reaction from the tower. Then the system may eventually cause significant dissipation of the kinetic energy if the tower is coupled to light fields in the thermal bath. To see what happens in the higher dimensional theory, we consider three models for the scenario, which are carefully prepared. In these models the extension is obvious but it does not disturb the original set-ups. The third model suggests that the evolution of the volume moduli may feel significant friction from the Kaluza-Klein tower.

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## I. INTRODUCTION AND THE MODEL

The study of the nonequilibrium dynamics, when a background field is time-dependent, can be motivated in many different physics of interest. Focusing on the cases of cosmology, one can find vast variety of applications in which cosmological events are studied using nonequilibrium dynamics. For instance in the inflation cosmology, nonequilibrium particle production has been studied for particle production after and during inflation [1–3], and for the dissipative process in warm inflation [3–6].

The nonequilibrium dynamics that leads to dissipative process during inflation can be studied assuming that the system is coupled to the thermal bath, or alternatively assuming physical decay factor in the propagator, to make the local approximation valid for the calculation. These studies have been developed using various techniques of the quantum field theory and the quantum statistical mechanics. A plausible scenario of dissipation has been studied in Ref. [4, 5] for the two-step model, using the methods of the quantum field theory when the local approximation can be applied. In this paper we will study the dissipation in the two-step model, when the Kaluza-Klein tower of the intermediate field changes its spectrum when a field moves.

Sometimes the dissipative equation of motion has been discussed specifically for the warm inflation scenario and thus there might be confusion between the criticism of the specific inflation scenario and the one for the calculation of the dissipative coefficients. In this paper we will not assume specific inflation background during dissipation, since the assumption of inflation requires additional (model-dependent) conditions for the system, to keep the thermal equilibrium a good approximation during the evolution. Indeed, dissipative equations may play significant role in non-inflating cosmological scenarios such as the curvaton [7] or the Affleck-Dine baryogenesis [8, 9]. The purpose of this paper is to show a clear example in which the Kaluza-Klein tower causes significant two-step dissipation in the low-temperature limit (i.e, when the temperature of the thermal bath is lower than the mass of the intermediate Kaluza-Klein mode.).

Model-dependent analyses are avoided in this paper. Previous studies on string-motivated models can be found in Ref.[10], in which the strategies are different from ours.

Let us start with the plain model. Consider the four-dimensional supersymmetric model with the interaction given by the superpotential [4, 5]:

$$W = W(\Phi) + g\Phi X^2 + hXY^2, \quad (1)$$

where  $\Phi$ ,  $X$  and  $Y$  are superfields that contains bosonic components  $\phi$ ,  $\chi$  and  $y$ , respectively [11]. In this model,  $X$  is the intermediate field that can decay into light field  $Y$ . The effective mass of the intermediate field is time-dependent because  $\Phi$  is moving. The time-dependent mass of the intermediate field causes particle creation and the dissipation of the kinetic energy. We are considering the effective equation of motion for the field  $\phi$ , which may contain the dissipative term when the local approximation is valid for the nonequilibrium dynamics. The superpotential  $W(\Phi)$  is not mandatory, since the potential could be obtained from a supersymmetry breaking term that lifts a flat direction of the supersymmetric theory. The auxiliary field equation of the field  $X$  gives

$$F_X^* + 2g\Phi X + hY^2 = 0, \quad (2)$$

which leads to interactions  $\sim (gh\phi)\chi y^2$  and  $\sim g^2\phi^2\chi^2$ . These terms are necessary for the later calculation of the dissipative coefficients.

Then, as far as the local approximation is valid, the dissipative term appears in the equation as

$$\ddot{\phi} + (3H + \gamma)\dot{\phi} + V_\phi = 0, \quad (3)$$

where the subscript of the scalar potential  $V(\phi)$  denotes the derivative with respect to the field. Here  $\gamma$  is the dissipative coefficient calculated in Ref.[5]. In the low-T regime ( $T/m_\chi \ll 1$ ), the coefficient is given by [4, 5]

$$\gamma \simeq g^2 h^4 \sum_{i=R,I} \left( \frac{g\phi}{m_{\chi_i}} \right)^4 \frac{T^3}{m_{\chi_i}^2}, \quad (4)$$

where  $m_{\chi_R}$  and  $m_{\chi_I}$  are the mass of the real and the imaginary parts of the bosonic component  $\chi$ . We will

have  $m_{\chi_R} \simeq m_{\chi_I}$  since the split is usually negligible. Here the interaction  $\sim hm\chi y^2$  ( $m \equiv g\phi$  in the above model) is crucial for the leading-order calculation<sup>1</sup>.

The model realizes the typical two-step dissipative process and avoids several difficulties of the warm inflation scenario. We thus consider the supersymmetric two-step scenario as a reference model. Inflation background is not necessary for the process.

In this paper we will consider the higher dimensional extension of the above two-step model. For the effective action that describes interaction with the Kaluza-Klein tower, we will consider the higher dimensional supersymmetry in the four-dimensional superspace [12], in which we can introduce a Gaussian distribution for the localization in the fifth dimension [13]:

$$f_G(y) = \frac{1}{\sqrt{2\pi}l_s} e^{-\frac{y^2}{2l_s^2}}, \quad (5)$$

where  $l_s^{-1}$  determines the width of the localization. The localization is needed to discuss the interaction between fields on the brane (boundary) and the fields in the bulk. The above Gaussian function is useful since it can explain the natural cut-off for the Kaluza-Klein tower [13]. However, just for the simplicity of the calculation, we will use  $\delta(y)$  instead and consider a simple cut-off at  $m_{KK} \simeq M_* \simeq l_s^{-1}$ . The replacement may affect numerical coefficients but does not change the qualitative result.

### A. Intermediate fields in the bulk

We consider a five-dimensional space time metric with the fifth dimension being compactified on a  $S^1/Z_2$  orbifold. The strategy is usually considered for obtaining four-dimensional MSSM from the flat five-dimensional action [12]. For the first model, we will place the intermediate field ( $X$ ) in the bulk while localizing  $\Phi$  and  $Y$  on the boundary.

From a four-dimensional perspective, the Kaluza-Klein towers of the intermediate field  $X$  can be rearranged in the form of  $\mathcal{N} = 2$  hypermultiplets. Choosing  $Z_2$  parity of the five-dimensional fields one can break the  $\mathcal{N} = 2$  supersymmetry to  $\mathcal{N} = 1$  supersymmetry. In  $\mathcal{N} = 1$  four-dimensional superspace, the five-dimensional hypermultiplet consists of the four-dimensional chiral superfields  $X(y)$  and  $X^c(y)$ , which are both labeled by  $y$  (the coordinate of the fifth dimension). The free action is given by [12]

$$S = \int d^5x \left\{ [\bar{X}^c X^c + \bar{X} X]_{\theta^2 \bar{\theta}^2} + (X^c \partial_y X|_{\theta^2} + \text{h.c.}) \right\}. \quad (6)$$

We will introduce the superpotential

$$W = [g\Phi X^2 + hXY^2] \delta(y), \quad (7)$$

which involves the bulk-boundary interactions. The auxiliary fields for the intermediate field  $X$  and  $X^c$  are computed to be:

$$\begin{aligned} F_X^* &= -2g\Phi X \delta(y) - hY^2 \delta(y) - \partial_y X^c \\ F_{X^c}^* &= \partial_y X. \end{aligned} \quad (8)$$

We assume the Fourier expansion on  $S^1/Z_2$ :

$$\begin{aligned} \chi &= \frac{1}{\sqrt{\pi R}} \chi^0 + \sqrt{\frac{2}{\pi R}} \sum_{n=1} \cos \left[ \frac{ny}{R} \right] \chi^{(n)} \\ \chi^c &= \sqrt{\frac{2}{\pi R}} \sum_{n=1} \sin \left[ \frac{ny}{R} \right] \chi^{c,(n)}, \end{aligned} \quad (9)$$

which suggests that the 0-mode appears only in  $\chi$ .

After integration, the effective four-dimensional superpotential for the KK state can be written in the form  $W = \sum_{m,n} W^{(m,n)}$ , where

$$W^{(m,n)} = g^{(m,n)} \Phi \left( X^{(m)} X^{(n)} \right)^2 + h^{(n)} X^{(n)} Y^2. \quad (10)$$

If the matrix “were” diagonal ( $g^{(m,n)} \simeq g^{(n)}$ ), the leading order calculation for the Kaluza-Klein tower ( $n \geq 1$ ) gives the dissipative coefficient

$$\gamma_{KK} \sim \sum_{n=1}^{n_{Max}} (g^{(n)})^2 (h^{(n)})^4 \left[ \frac{g^{(n)} \phi}{m_{\chi}^{(n)}} \right]^4 \frac{T^3}{[m_{\chi}^{(n)}]^2}, \quad (11)$$

where  $n_{Max} \simeq RM_*$  gives the cut-off at  $m_{KK} \sim M_*$ . The off-diagonal elements may contribute to the multiplicity of the diagram. Let us examine the corrections from the off-diagonal elements.

If one assumes <sup>2</sup>  $g\phi \ll 1/R$ , the mass eigenstates can be approximated by the pure KK modes and then one immediately finds

$$\begin{aligned} \gamma_{KK} &\sim \sum_{n=1}^{n_{Max}} \sum_{m=1}^{n_{Max}} \frac{(g^{(m,n)})^6 (h^{(n)})^2 (h^{(m)})^2 T^3}{(m_{KK}^{(n)})^3 (m_{KK}^{(m)})^3} \phi^4 \\ &\sim g^6 h^4 T^3 R^6 \phi^4 \sum_{n=1}^{n_{Max}} \sum_{m=1}^{n_{Max}} \frac{1}{n^3 m^3} \\ &\sim g^6 h^4 T^3 R^6 \phi^4 \\ &< \gamma_0, \end{aligned} \quad (12)$$

<sup>1</sup> See Ref.[5] for the details of the calculation.

<sup>2</sup> This condition is not mandatory since the extra dimension could be large ( $R > M_*^{-1}$ ). This possibility will be considered later in this section.

where  $\gamma_0 \equiv C_0 h^4 \frac{T^3}{\phi^2}$  is the dissipative coefficient for the intermediate 0-mode. Therefore, in the above limit ( $g\phi \ll 1/R$ ) the Kaluza-Klein tower gives negligible contribution to the dissipative coefficient. As a result, we find

$$\gamma \simeq C_0 h^4 \frac{T^3}{\phi^2} + C_{KK} g^6 h^4 T^3 R^4 \phi^2 \sim \gamma_0, \quad (13)$$

where  $C_0$  and  $C_{KK}$  are numerical coefficients.

Significant enhancement is possible in the opposite limit, when  $g\phi \gg R^{-1}$ . However, replacing the Gaussian function by  $\delta(y)$ , one will find off-diagonal elements of  $g^{(m,n)}$ , which are exactly the same size as the diagonal ones. For later use, we first estimate the contribution from the diagonal elements. Here we consider

$$W^{(n)} = g^{(n)} \Phi \left( X^{(n)} \right)^2 + h^{(n)} X^{(n)} Y^2. \quad (14)$$

The leading-order contribution gives

$$\begin{aligned} \gamma_{KK} &\sim n_* h^4 \frac{T^3}{\phi^2} \\ &\sim g R h^4 \frac{T^3}{\phi}, \end{aligned} \quad (15)$$

where  $n_* \sim gR\phi$  is not the usual cut-off ( $n_{Max}$ ) of the Kaluza-Klein tower, but obtained from the condition  $m_{KK} < g\phi$ . The previous calculation for  $g\phi \ll m_{KK}^{(1)}$  suggests that the Kaluza-Klein states in  $n \geq n_*$  is negligible and does not contribute to the dissipation. Finally, we find

$$\gamma \sim C_0 h^4 \frac{T^3}{\phi^2} + \hat{C}_{KK} g R h^4 \frac{T^3}{\phi}. \quad (16)$$

The above result suggests that there could be an interesting critical point  $\phi_c$ , where the dissipation becomes strong due to the Kaluza-Klein tower. We can calculate  $\phi_c$  as

$$\phi_c \simeq \frac{C_0}{g \hat{C}_{KK}} R^{-1}. \quad (17)$$

What is the consequence of  $\phi_c$ ? If  $\phi$  is rolling down on a hilltop potential, it may feel significant friction when  $\phi$  is small<sup>3</sup>. In that case the friction-dominated motion can be prolonged by the Kaluza-Klein states. In that way the hilltop inflation [16] (or the hilltop curvaton [17]) could be affected by the Kaluza-Klein tower. On the other hand, if the field is moving toward the origin, the Kaluza-Klein enhancement ceases toward  $\phi_c$ . This may affect the starting point of the curvaton oscillation and/or the Affleck-Dine baryogenesis after inflation [7–9].

In the above scenario, the field  $\phi$  is localized on the boundary, which means that the vacuum expectation value of the field is bounded from above ( $\phi \leq M_*$ ). Therefore, the significant Kaluza-Klein enhancement is possible only when the extra dimension is large ( $RM_* \gg 1$ ) and  $\phi$  is moving in the region  $\phi_c < \phi < M_*$ .

Here, our concern is the contributions from the off-diagonal elements. To understand the physics related to the off-diagonal elements, let us consider 2-field model with the superpotential

$$W = g\Phi(X_1 + X_2)^2 + h(X_1 + X_2)Y^2. \quad (18)$$

Then one can choose  $X_{\pm} \equiv (X_1 \pm X_2)/\sqrt{2}$  to obtain

$$W = (2g)\Phi X_+^2 + (\sqrt{2}h)X_+ Y^2. \quad (19)$$

For the field  $X_+$ , the dissipation coefficient is multiplied by  $(\sqrt{2}h/h)^4$ . Applying the same multiplicity, we find

$$\begin{aligned} \gamma_{KK} &\sim n_*^2 h^4 \frac{T^3}{\phi^2} \\ &\sim (gR)^2 h^4 T^3. \end{aligned} \quad (20)$$

Note however the weak-coupling expansion is not valid when  $g_{\text{eff}} \equiv n_* g \geq 1$  or  $h_{\text{eff}} \equiv \sqrt{n_*} h \geq 1$ . Therefore, a modest argument would be that Eq.(20) certifies significant dissipation, but the higher order contributions could be larger than the above result. At this moment we do not calculate the higher contributions in the strong limit because it is not suitable for the purpose of this paper.

## B. $\Phi$ and $X$ in the bulk

Second, assume that  $\Phi$  and  $X$  are both in the bulk while  $Y$  is localized on the boundary [15]. We will introduce the superpotential for the interaction

$$W = [g\Phi X^2 + hXY^2] \delta(y), \quad (21)$$

where the “point interaction” is assumed for the Yukawa coupling  $g\Phi X^2$ . This assumption could be removed, however the conventional Yukawa interaction is forbidden in the  $\mathcal{N} = 2$  supersymmetry. The effective four-dimensional superpotential that contains the 0-mode  $\Phi^{(0)}$  and the Kaluza-Klein state is thus given by

$$W^{(m,n)} = g^{(m,n)} \Phi^{(0)} \left( X^{(m)} X^{(n)} \right)^2 + h^{(n)} X^{(n)} Y^2. \quad (22)$$

The above superpotential involves the interactions that are required for the two-step dissipation. Only the 0-mode of  $\phi$  is moving in that model. All the other fields are supposed to have negligible vacuum expectation values.

The significant difference appears not in the “calculation” but in the “condition” that constrains the maximum value of  $\phi$ . In contrast to the previous model, in

<sup>3</sup> However,  $\phi$  must be larger than the temperature of the thermal bath so that the low-temperature approximation is valid for the calculation.

which  $\Phi$  is localized, one is not forced to put the condition  $\phi \leq M_*$ . Therefore, for the diagonal elements we find  $\gamma_{KK} \sim \hat{C}_{KK} g R h^4 \frac{T^3}{\phi}$ . Again, including the possible multiplicity from the off-diagonal elements, we find  $\gamma_{KK} \sim \check{C}_{KK} (gR)^2 h^4 T^3$ , where  $\check{C}_{KK}$  is the numerical coefficient.

### C. Dissipation when radius changes

Note first that  $m_{KK}^{(n)}$  is basically the function of  $R$ , which could be a dynamical parameter in the early Universe. Then a natural question arises: what is the consequence of the Kaluza-Klein tower that is changing its spectrum due to the dynamical moduli?

In this section we will define the moduli field  $\Phi \equiv M_*^2 R$  and consider the motion of  $\Phi$ . For simplicity, the canonical kinetic term is assumed for the field  $\Phi$ , however this assumption is not mandatory.  $X^c$  and  $X$  are in the bulk as usual, and  $Y$  is localized on the boundary. Consider the Sherk-Schwartz (SS) boundary condition [14] that replaces  $n \rightarrow n + \omega$  in the Fourier expansion of the bosonic field. We will introduce the interaction

$$W = h X^c Y^2 \delta(y - \pi R), \quad (23)$$

which leads to the auxiliary field equation

$$F_{X^c}^* = -h Y^2 \delta(y - \pi R) + \partial_y X. \quad (24)$$

The superpotential gives the interaction required for the two-step dissipation at the leading order. The crucial interaction is

$$\mathcal{L}_{\text{int}} \sim h^{(n)} m_{KK}^{(n)} \chi^{(n)} y^2, \quad (25)$$

where  $m_{KK}^{(n)}(\Phi) \simeq n M_*^2 / \Phi$  is the Kaluza-Klein mass expressed using  $\Phi$ . Again, we find the dissipative coefficient

$$\begin{aligned} \gamma_{KK} &\sim \sum_{n=1}^{n_{Max}} \frac{(m_{KK}^{(n)})^2}{\Phi^2} h^4 \frac{T^3}{(m_{KK}^{(n)})^2} \\ &\sim h^4 \frac{T^3}{M_* \Phi}, \end{aligned} \quad (26)$$

where  $n_{Max} \sim R M_* \sim \Phi / M_*$  is used in the last line.

Note that the spectrum of the tower is changing whenever the moduli is time-dependent. There is no need to consider additional interaction between  $\Phi$  and  $X$ .

The above model might be a reminiscence of Casimir force [18, 19]. However, in the present model the creation of the light field causes irreversible process. They must be discriminated.

## II. CONCLUSIONS AND DISCUSSIONS

In this paper we considered two-step dissipation for the supersymmetric model in which the intermediate field  $X$  is placed in the bulk of the fifth dimension. In the effective action, the intermediate field has the tower of Kaluza-Klein states that are coupled to the light fields on the boundary.

We first considered a simple generalization of Ref.[5]. We first estimated the contribution from the diagonal interactions. The Kaluza-Klein tower of the intermediate field  $X$  enhances the dissipative coefficient when  $\phi > \phi_c \equiv \frac{C_0}{g \check{C}_{KK}} R^{-1}$ . In that way the dissipative coefficient changes from  $\gamma_0 = C_0 h^4 \frac{T^3}{\phi^2}$  to  $\gamma \sim C_0 h^4 \frac{T^3}{\phi^2} + \hat{C}_{KK} g R h^4 \frac{T^3}{\phi}$  for the diagonal elements, and to  $\gamma \sim C_0 h^4 \frac{T^3}{\phi^2} + \check{C}_{KK} (gR)^2 h^4 T^3$  when the multiplicity from the off-diagonal elements are included. More enhancement is expected in the strong ( $g_{\text{eff}}, h_{\text{eff}} > 1$ ) limit.

Then we considered a model with a moving moduli, in which the moduli determines the mass of the Kaluza-Klein states. The result suggests that the back-reaction from the Kaluza-Klein tower could be crucial for the dynamics of the expanding extra dimension.

We have some comments about the simplification. In this paper we have disregarded the Kaluza-Klein graviton, which might be important for the dynamics. We have also disregarded the decay from the higher to the lower Kaluza-Klein states. Both of those may enhance the dissipative process.

Warped extra dimension has not been considered in this paper. In that model the fields on the brane can change their mass when the brane moves along a throat.

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